

COLLAPSE OF FLUX TUBES

L. Wilets and R. D. Puff

Department of Physics, University of Washington, Seattle, WA 98195, USA

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Abstract

The dynamics of an idealized, infinite, MIT-type flux tube is followed in time as the interior evolves from a pure gluon field to a $\bar{q} q$ plasma. We work in color U(1). $\bar{q} q$ pair formation is evaluated according to the Schwinger mechanism using the results of Brink and Pavel. The motion of the quarks toward the tube endcaps is calculated by a Boltzmann equation including collisions. The tube undergoes damped radial oscillations until the electric field settles down to zero. The electric field stabilizes the tube against pinch instabilities; when the field vanishes, the tube disintegrates into mesons. There is only one free parameter in the problem, namely the initial flux tube radius, to which the results are very sensitive. Among various quantities calculated is the mean energy of the emitted pions.

I. INTRODUCTION

Flux tubes are one of the most elementary systems of quantum chromodynamics. They are the idealized configurations of heavy quark-antiquark pairs at large separation L such that the region between can be assumed to possess axial-cylindrical symmetry. They play a central role in lattice QCD calculations and in models of QCD, as well as in the phenomenology of QCD processes, such as jets produced in heavy ion collisions or the “spaghetti” connecting the separating heavy ions.

From the spectroscopy of heavy quarkonia and Regge trajectories, it has been possible to extract a reliable measurement of the string tension by assuming the two body potential for heavy quarks to be of the form

$$V(L) = -\frac{4\alpha_s}{3L} + \theta L + \text{constant} \quad (1)$$

The “experimental” value is [1] $\theta = 913 \text{ MeV/fm} = 0.18 \text{ GeV}^2 = 4.63/\text{fm}^2$. In fact the linear region of the flux tube potential is not directly explored by quarkonia or by lattice QCD calculations:

The wave functions in quarkonia calculations tend to be centered near the “knee” of the potential although high-lying states do explore more deeply into the linear region. However other forms of the heavy quark potential (including a logarithmic term) have also achieved some success in reproducing quarkonia spectra.

Lattice QCD calculations on flux tubes are generally limited to the quenched approximation (no massless quarks) and allow for a separation of the heavy quark-antiquark of only about 1 fm, which do not reproduce a cylindrical region

Static flux tubes are unstable at separations greater than 1 fm, since the energy required to stretch the tube by 1 fm is nearly 1 GeV, which is about the energy difference between a quarkonium, $Q\bar{Q}$, and a pair of heavy-light mesons, $Q\bar{q} + \bar{Q}q$. Lattice calculations without light quarks cannot explore this instability.

Recently there have been several papers [2–5] which have explored the creation of light quark-antiquark pairs as a mechanism for flux tube breaking. They employ a variation of the Schwinger [6] parallel plate capacitor model in which the infinite transverse geometry is replaced by the MIT boundary condition

$$-i\boldsymbol{\gamma}\cdot\boldsymbol{\rho}\psi(R) = \psi(R) \quad (2)$$

at the cylindrical tube surface for the light quarks. The electric field \mathbf{E} is constant over all space in the longitudinal (z) direction, with \mathbf{D} equal to \mathbf{E} inside the tube and zero outside. It is found that pair production in the cylinder is suppressed relative to the Schwinger formula due to transverse confinement which gives the quarks an effective mass proportional to $1/a$ where a is the radius of the flux tube.

In the present work, we make some simplifying assumptions, the most significant are:

We consider infinite flux tubes. This means that the tube is long compared with its radius. For a jet, this means that the characteristic tube lifetime, $\tau \gg a/c$, since the quark producing the jet travels near the speed of light.

We use the MIT bag model for the tubes.

Although some contact is made with color SU(3), we perform our calculations in U(1). The extension to Abelian SU(3) is not a serious complication.

We adopt one dimensional Debye screening for the quark-quark interactions used in the Boltzmann collision term. This is further simplified to an effective δ -function potential.

With these approximations, our results depend on only one parameter, the initial (static) flux tube radius.

II. THE MIT MODEL

We utilize the framework of the MIT model because it is so very simple. It is the prototype of color dielectric models. Some features of various models of flux tube are discussed in Sec. VIII.

In the MIT model, the calculation of the flux tube configuration is straightforward since it involves no quarks. The flux through the tube is equal to the color charge $Q = \frac{1}{2}\boldsymbol{\lambda}g_s$, where the $\boldsymbol{\lambda}$ are the Gell-mann color matrices. The electric field is $E = Q/A$. Here Q is the charge at one end (say, the left) of the flux tube (\bar{Q} at the other). It is this charge which will be shielded as a function of time due to the uniform production of pairs out of the sea in the flux tube. We distinguish it from the charge q of the light quarks being created in pairs, although for the static model $Q = q$.

The energy per unit length is given by

$$\frac{\mathcal{E}}{L} = B A + \frac{1}{2} E^2 A = B A + \frac{Q^2}{2A} \quad (3)$$

where $A = \pi a^2$, with a the cylinder radius, and B is the bag constant. Minimization with respect to A yields $A = Q/\sqrt{2B}$ and

$$\left. \frac{\mathcal{E}}{L} \right|_{min} \equiv \theta = \left[2 Q^2 B \right]^{1/2} = 2 B A = \frac{Q^2}{A}. \quad (4)$$

The translation to QCD obtains by the replacement of Q^2 by its expectation value in a color-singlet state,

$$Q^2 \rightarrow \langle Q^2 \rangle = \frac{g_s^2}{4} \langle \lambda^2 \rangle = \frac{16 \pi \alpha_s}{3}, \quad (5)$$

where we have used $\langle \lambda^2 \rangle = 16/3$ and $\alpha_s = g_s^2/4\pi$.

III. THE SCHWINGER MECHANISM

The rate of pair creation in the Schwinger mechanism with MIT boundary conditions has been calculated by several groups [2], [3], [4]. The point of departure is the QED Schwinger problem of infinite parallel plate capacitors infinitely far apart [6]. The rate of pair production per unit time per unit volume is

$$w(\infty) = \frac{\kappa^2}{4\pi^3} \sum_{n=1}^{\infty} \frac{e^{-\pi n m_0^2/\kappa}}{n^2} \rightarrow \frac{\kappa^2}{4\pi^3} \frac{\pi^2}{6} = \frac{\kappa^2}{24\pi} \quad \text{for } m_0 \rightarrow 0. \quad (6)$$

where m_0 is the Fermion mass and $\kappa = |q \mathbf{E}|$.

In the MIT model, κ (with $q = Q = \sqrt{\theta\pi}a$) is just the string tension θ . in which case

$$w(\infty) \approx 0.25 \text{ c/fm}^4 \quad (7)$$

This identification is model-dependent, as discussed below. For example, Pavel and Brink [2] identify $\kappa = 2\theta$. (For reference, we note that those authors denote the string tension by σ whereas we use θ .)

As we will see below, the rate is strongly suppressed by confinement.

IV. CONFINEMENT AND TRANSVERSE MASS

Transverse confinement has the effect of introducing a “transverse mass” in the formula for w : $m_0^2 \rightarrow m_0^2 + x_n^2/a^2$. Using the results of Pavel and Brink [2] [*cf.* their Eq. (4.3)], the pair production rate per unit time per unit length W is

$$W(a) = -\frac{n_f \kappa}{\pi} \sum_{x_n > 0} \ln \left[1 - e^{-\pi x_n^2/a^2 \kappa} \right] \approx \frac{n_f \kappa}{\pi} e^{-\pi x_1^2/a^2 \kappa}, \quad (8)$$

where n_f is the number of flavors. For reasonable a and $m_0 = 0$, the two lowest transverse modes, $x_{\pm 1}$ dominate, and we consider these only here: $|x_{\pm 1}| = 1.4347$. Henceforth we will assume $m_0 = 0$.

V. SIMPLE ESTIMATE OF COLLAPSE TIME

In the following, we work in QED, which is to say, charge is a scalar quantity, plus or minus.

We now consider $Q = Q(t)$, since the charge at the ends is shielded due to the flow of opposite charge into the ends. We will calculate it more carefully in the next section, but it is a good approximation to calculate the current by using the initial radius of the flux tube and the electric field for short times. The number of pairs is proportional to the time, and the mean velocity of a quark is proportional to the time. It is then easy to show that

$$\frac{dQ(t)}{dt} = -2 J \approx -2 W(0) q \frac{\langle p \rangle}{m} t \approx -2 \frac{W(0) q^2}{2 A m} t^2 Q(0), \quad (9)$$

where J is the current carried by positive quarks; negative quarks carry an equal current, including the same sign. Hence the factor of 2. We have used the average velocity of quarks $\langle p \rangle / m \approx Q(0) q t / (2 A m)$, We make the identification $m = m_{\perp} = x_1/a$. Again we note that a is the tube radius at $t = 0$. One more time integration gives

$$Q(t) = Q(0) \left[1 - \frac{W q^2}{3 A m} t^3 \right]. \quad (10)$$

We now estimate the characteristic lifetime of the tube to be

$$\tau(a) \approx \left[\frac{3 m A}{W q^2} \right]^{1/3} = \left[\frac{3 m}{W \kappa} \right]^{1/3} = \left[\frac{3 \pi x e^{\pi x_1^2/a^2 \kappa}}{\kappa^2 a} \right]^{1/3}. \quad (11)$$

Note that there is no dependence on the length of the flux tube!

Taking $n_f = 2$, we find

$$\kappa = \theta \quad 2\theta \quad (12)$$

$$\tau(0.5) = 7.0 \quad 1.7 \quad \text{fm}/c, \quad (13)$$

$$\tau(1.0) = 1.4 \quad 0.68 \quad \text{fm}/c. \quad (14)$$

Recall that $\kappa = \theta$ in the MIT model.

VI. COLLISIONLESS BOLTZMANN EQUATION

Here we still employ the adiabatic MIT model to follow the dynamics of the flux tube wall by assuming that the tube radius is in instantaneous equilibrium. In addition to the electric field, the light quark plasma also exerts a pressure on the tube wall. We assume that only the lowest transverse eigenmode is excited ($x_{\pm 1}$), so that the tube radius, now denoted by $R(t)$ instead of a , is determined by minimizing the transverse energy

$$\frac{\mathcal{E}}{L} = \pi B R^2 + \frac{Q^2}{2\pi R^2} + \frac{2 N x_1}{R}, \quad (15)$$

which yields

$$-\frac{Q^2}{\pi} + 2\pi B R^4 - 2 N x_1 R = 0. \quad (16)$$

$N = N(t)$ is the number of *pairs* per unit length,

$$N(t) = \int_{-\infty}^{\infty} n(p, t) dp = \int_0^t W(t') dt \quad (17)$$

with $n(p, t)$ the momentum density per unit length for either q or \bar{q} . Note that we distinguish between the shielded heavy quark charge $Q(t)$ and the elementary light quark charge q , where $Q(0) = q$.

The pair production rate (per unit length) is given by

$$W = n_f \frac{\kappa}{\pi} e^{-\pi x_1^2 / R^2 \kappa}, \quad (18)$$

with $\kappa = q|E|$.

The transverse motion of the quarks is now treated relativistically but classically. Thus the relativistic mass is taken to be

$$m^2(t) = \left(\frac{x_1}{R(t)} \right)^2 + p^2. \quad (19)$$

The rate of change of $n(p, t)$ is governed by the coupled integro-differential equations:

$$\frac{\partial n}{\partial t} = -q E \frac{\partial n}{\partial p} + \delta(p) W + \left(\frac{dn}{dt} \right)_{col}, \quad (20)$$

$$J(t) = q \int dp n p / m, \quad (21)$$

$$\frac{dQ}{dt} = -2 J, \quad (22)$$

$$E(t) = \frac{Q(t)}{\pi R(t)^2}, \quad (23)$$

where $(dn/dt)_{col}$ is the rate of change of the momentum distribution due to collisions. The initial conditions are $n(p, 0) = 0$, $Q(0) = q = \sqrt{\theta\pi}a$, $R(0) = a$.

We wish to point out the very interesting fact that **no magnetic fields** are produced! This is because, in Maxwell's famous fourth equation, the convection current is exactly cancelled by the displacement current.

One valuable check on the numerics is provided by the agreement between the two forms of Eq. 17. Excellent agreement is attained until numerical instability sets in, when the two values diverge dramatically, in which case time or momentum steps were changed until equality was achieved.

These equations have been solved numerically first without the collision term for several parameter sets. The results are displayed in figure 2; the thin lines are for the collisionless solutions. We note here oscillations with negligible damping. In the absence of damping, the current continues to flow even when the charge at the end caps (and hence \mathbf{E} goes through zero. Consequently the charge at the end caps changes sign, tending to reverse the current. The number of pairs increases with time. The energy per unit length also increases; there is no energy conservation within a segment of the tube.

VII. COLLISIONAL THERMALIZATION AND DAMPING

The problem is simplified by the one-dimensional nature of the dynamics: we assume that the transverse motion can be handled by the adiabatic MIT model. Scattering between quarks (charge q) or between quarks and anti-quarks (charge $-q$) is like that of beads on a string. The colliding particles exchange momentum with some probability to be computed below. Collisions between identical particles does not alter the momentum distribution – it is as though no collision took place. Since the interaction is independent of spin and flavor, we do not consider any quark-quark collisions, but only collisions between quarks and antiquarks. Then the change in population density of quarks in momentum satisfies the equation

$$\begin{aligned} \left(\frac{dn_+(p_1)}{dt} \right)_{col} &= \int dp_2 |\mathcal{R}(p_1 - p_2)|^2 (v_1 - v_2) [n_-(p_1)n_+(p_2) - n_+(p_1)n_-(p_2)] \\ &= \int dp_2 |\mathcal{R}(p_1 - p_2)|^2 (v_1 - v_2) [n(-p_1)n(p_2) - n(p_1)n(-p_2)] \end{aligned} \quad (24)$$

with $n = n_+$. Here $|\mathcal{R}|^2$ is the reflection probability for the collision of quarks with antiquarks in a one dimensional potential barrier or well. The potential should take into account shielding introduced by the presence of other quarks and antiquarks.

A. The Shielded Quark-Antiquark Interaction

The interaction between confined quarks is that of disks in the tube of radius R . The confinement of the D -field by the tube walls leads to a potential between bare quarks which is linear in the separation, just like the flux tube potential itself. However the quarks are shielded by other quarks and antiquarks. An estimate of the shielding can be obtained from a one-dimensional version of the Debye formula:

The charge density in the vicinity of a quark at (say) $z = 0$ in a thermal bath in the rest frame of the quark is

$$\rho(z) = \frac{q}{\pi R^2} \delta(z) + C \int_{-\infty}^{\infty} dp \left[e^{\epsilon^+(p)/k_B T} - e^{\epsilon^-(p)/k_B T} \right] \quad (25)$$

where

$$\epsilon^\pm(p) = \pm q \phi(z) + \sqrt{m_0^2 + p^2} \quad (26)$$

with $m_0 = x_1/R$ and $\phi(z)$ the shielded potential to be determined. Since $\phi(z = \pm\infty) = 0$, it follows that the polarization charge of $+/-$ quarks is

$$\rho^\pm(x = \pm\infty) = \pm \frac{q N}{\pi R^2} = \pm C \int_{-\infty}^{\infty} e^{-\sqrt{m_0^2 + p^2}/k_B T} dp \quad (27)$$

and hence

$$\rho(z) = \frac{q}{\pi R^2} \delta(z) + \frac{q N}{\pi R^2} \left[e^{-q\phi(z)/k_B T} - e^{q\phi(z)/k_B T} \right] \quad (28)$$

If, as turns out to be the case generally, $|q\phi|/k_B T \ll 1$, the expression can be linearized in ϕ , and we obtain

$$\rho(z) = \frac{q}{\pi R^2} \delta(z) - \frac{2N q^2}{\pi R^2 k_B T} \phi(z) \quad (29)$$

Poisson's equation in one dimension then reads

$$\frac{d^2 \phi}{dz^2} = -\rho(z) = -\frac{q}{\pi R^2} \delta(z) + \frac{2N q^2}{\pi R^2 k_B T} \phi(z) \quad (30)$$

The solution to this equation is given by

$$\phi(z) = \frac{q \lambda_D}{2\pi R^2} e^{-|z|/\lambda_D} \quad (31)$$

with the Debye length given by

$$\lambda_D^2 = \frac{\pi R^2 k_B T}{2 N q^2} \quad (32)$$

Therefore the quark-antiquark potential is given by

$$V_D(z) = -q\phi(z) = -\frac{q^2 \lambda_D}{2\pi R^2} e^{-|z|/\lambda_D} \quad (33)$$

B. Calculation of the reflection probability

It is an elementary quantum mechanical problem to calculate $|\mathcal{R}|^2$ for the one dimensional Debye potential Eq. 33. For greater simplicity and insight, we consider two analytic approximations for $|\mathcal{R}|^2$:

(1) A square-well potential of width D and depth V_0 such that $V_0 D$ equals the integral of $V_D(z)$. A reasonable value for D is $2\lambda_D$. The reflection probability can be found in any of several texts.

(2) A delta-function potential of strength V_δ (like the square-well potential with $D \rightarrow 0$). The strength is given by

$$V_\delta = -\frac{k_B T}{2N} \quad (34)$$

C. Evaluating the temperature

In order to identify the temperature T , we note that, except for very small times, the charactersitic momenta in the problem are large compared with m , and therefore we may set $\epsilon^\pm \approx \pm q\phi + |p|$ in Eq. 26. Consequently we can approximate [cf. Eq. 25]

$$n(p) \equiv n_+(p) \approx C e^{-|p - \langle p \rangle|/k_B T}, \quad (35)$$

where $\langle p \rangle$ is the average drift momentum. For this form of the Boltzmann distribution, we have

$$N = \int_{-\infty}^{\infty} n(p) dp = 2 C T \quad (36)$$

and the fluctuation in the momentum yields the temperature:

$$(k_B T)^2 = (2N)^{-1} \int_{-\infty}^{\infty} n(p) [p - \langle p \rangle]^2 dp \quad (37)$$

In the calculation, we use these equations to determine the temperature at each time from the calculated $n(p, t)$. Then

$$(k_B T)^2 = \langle (p - \langle p \rangle)^2 \rangle / 2 \quad (38)$$

where $\langle \rangle$ means momentum average with respect to the calculated $n(p, t)$.

D. \mathcal{R} for the delta-function

The delta-function case $V(z_1 - z_2) = V_\delta \delta(z_1 - z_2)$ has a particularly simple form for the reflection amplitude \mathcal{R} and probability $|\mathcal{R}|^2$ for particles of momenta p_1 and p_2 , namely

$$|\mathcal{R}|^2 = \frac{V_\delta^2}{V_\delta^2 + (p_1 - p_2)^2 / \mu^2} \quad (39)$$

where, without rigor, we have used a reduced mass

$$\frac{1}{\mu} = \frac{1}{m(p_1)} + \frac{1}{m(p_2)}. \quad (40)$$

With the self-consistent determination of V_δ from Eqs. 39 and 40, a complete solution of Eq. 20 and comparison with the collisionless results for the same model parameter (only a) is possible.

E. Damping time

We see here the evolution from a *flux* tube to a *plasma* tube. The behaviour of this evolution depends strongly on the single parameter a . For a between 0.5 and 1.5 fm, we note damped oscillations which increase in amplitude and decreases in period as a increases. We proffer three measure of the characteristic damping time, τ_0 , τ_1 , and τ_2 . τ_0 is defined as the time at which $Q(t)$ first crosses zero. The other definitions are

$$\tau_1 \equiv \int_0^\infty |Q(t)| t dt \Big/ \int_0^\infty |Q(t)| dt \quad (41)$$

and

$$\tau_2 \equiv \int_0^\infty Q^2(t) t dt \Big/ \int_0^\infty Q(t)^2 dt \quad (42)$$

Calculated results for these three measures are shown in Fig. 6.

In Fig. 5 we show the number of pairs produced per unit length, and the energy per pion. The pion production is calculated assuming that one meson, π , ρ , or ω , is produced in proportion to their spin-isospin statistical weights: 3/15 for pions, 9/15 for rhos, and 3/15 for omegas. The rho and omega decay to 2 and three pions respectively. The total energy per unit length is distributed among the pions per unit length. The result is shown in Fig. 6.

Note that our calculations do not depend upon the motion of the end caps, and therefore, within our approximations, the energy per ion is flat in rapidity.

Our results show a high degree of sensitivity to the initial flux tube radius a . This provides a useful handle for distinguishing among various flux tube models and, in particular, the appropriate size of a .

F. Ultimate disintegration of the plasma tube

Unlike the plasmas in a fusion reactor, the *flux* tube is stable against the sausage instability; the difference is due to the absence of magnetic fields. After the electric field had decayed, this stabilizing effect is removed and the *plasma* tube falls apart. We have not attempted yet to follow this final phase, but use the various τ 's given above to estimate the disintegration time. Our own bias is to use τ_1 .

VIII. MODELS OF FLUX TUBES

Since the collapse time is strongly dependent upon the initial flux tube radius $R(0) = a$, we now address more realistic flux tube calculations to determine a and the relationship between κ and θ .

A. The MIT Model

In this model, the field energy $E^2/2$ is uniform within a radius a which is characteristically 1 fm or larger. The flux tube energy is shared equally between the electric field and volume energies: $\theta = E^2 A = Q^2/A = q|E|$.

B. Lattice Gauge Calculations

Lattice calculations at present are usually restricted to the quenched approximation (no zero-mass dynamical quarks) and to $Q\bar{Q}$ separations L of less than about 1 fm. Nevertheless Sommer [7] obtains a dependence of the (longitudinal) electric field energy which appears to be stable with L . The functional form is roughly exponential, $E^2/2 \propto e^{-r_\perp^2/b}$ with $\langle r_\perp^2 \rangle = 6b^2 \approx (0.2 \text{ fm})^2$. The *equivalent* radius (the radius of a square form giving the same $\langle r_\perp^2 \rangle$) is $r_{eq} = \sqrt{2} \langle r_\perp^2 \rangle^{1/2} \approx 0.3 \text{ fm}$. Here $\theta = 2q|E|$.

C. Dual Superconductivity Model

The dual superconductivity model [8] yields a somewhat larger radius than lattice calculations, with $r_{eq} \approx 0.4$ fm.

D. The CDM

The chromodielectric model [9] is similar to the MIT model, except that the electric field energy has a diffuse surface and the flux tube energy θ also contains a surface term such that $\theta > q|E|$.

IX. FURTHER CORRECTIONS

There are several effects which are yet to be included in order to improve the estimate of the collapse time:

(1) Finite length effects also increase the collapse time. There are at least two such effects to consider: a) The current associated with a given momentum group terminates when all the charge in the group is transported the full length of the tube. b) No pairs can be created unless the change in potential the length of the tube is greater than m_{\perp} : $|q E L| > x_1/R$. These also lengthens the collapse time.

(2) Dynamics of the confinement mechanism and associated energy conservation. One might consider the chromodielectric soliton model to handle this.

X. FINAL COMMENTS

It is important to note that the calculations presented here do not depend on the velocity of the heavy quarks in the longitudinal direction. The collapse time depends sensitively on the radius of the flux tube and indeed on detailed structure of the tube.

For different approaches to the problem, see Kluger, *et al.* [10] and Rau [11].

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FIGURES

FIG. 1. Schematic evolution of a flux tube to a plasma tube to a collection of mesons. The radius of the tube exhibits damped oscillations.

FIG. 2. Variation with time of the flux tube radius $R(t)/a$ (upper curves) and net endcap charge $Q(t)/q$ (lower curves) for various a . For the case (a) of $a = 0.5$ fm. we have also shown in thin lines the solution to the collisionless Boltzmann equation.

FIG. 3. The mean longitudinal momentum per quark as a function of time for $a = 0.5$ and 1.0 fm.

FIG. 4. The quark “temperature” as a function of time for $a = 0.5$ and 1.0 fm.

FIG. 5. The growth of the number of pairs as a function of time for $a = 0.5$.

FIG. 6. The variation of several important quantities with a : The characteristic damping times (in fm/c) τ_0 (where $Q(t)$ first crosses zero), τ_1 and τ_2 (as defined by equations 41 and 42). Ten times the number of pairs produced per fm is denoted by $10 N$. E is the final energy per fm ($1/\text{fm}^2$). E_π is the mean energy (in GeV) of pions produced by disintegration of the tube.

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